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On the BPS Spectrum at the Root of the Higgs Branch

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Abstract

We study the BPS spectrum and walls of marginal stability of the $\mathcal{N} = 2$ supersymmetric theory in four dimensions with gauge group $SU(n)$ and $n \leq N_f < 2n$ fundamental flavours at the root of the Higgs branch. The strong-coupling spectrum of this theory was conjectured in hep-th/9902134 to coincide with that of the two-dimensional supersymmetric \mathbb{CP}^{2n-N_f-1} sigma model. Using the Kontsevich–Soibelman wall-crossing formula, we start with the conjectured strong-coupling spectrum and extrapolate it to all other regions of the moduli space. In the weak-coupling regime, our results precisely agree with the semiclassical analysis of hep-th/9902134: in addition to the usual dyons, quarks, and W bosons, if the complex masses obey a particular inequality, the resulting weak-coupling spectrum includes a tower of bound states consisting of a dyon and one or more quarks. In the special case of \mathbb{Z}_n -symmetric masses, there are bound states with one quark for odd n and no bound states for even n .

1 Introduction

We consider the BPS spectrum of the $\mathcal{N} = 2$ SQCD with gauge group $SU(n)$ and $n \leq N_f < 2n$ fundamental flavours at the root of the Higgs branch in four dimensions. The central charge in this theory was shown in [2] to be the same as in the $\mathcal{N} = (2, 2)$ supersymmetric \mathbb{CP}^{2n-N_f-1} sigma model in two dimensions with twisted mass terms (with an appropriate identification of parameters). On this basis, the BPS spectra of the two theories were also conjectured to coincide¹. Additional support for the conjecture was provided [3], and an explanation for the coincidence of the spectra in terms of vortex strings was given in [9, 8]. In this paper, we present further evidence for this conjecture using the Kontsevich–Soibelman wall-crossing formula for the four-dimensional theory. The finite set of BPS states of the two-dimensional theory in the strong-coupling regime is known [1, 12]. Assuming that the strong-coupling spectrum of the four-dimensional theory is indeed the same, we find the walls of marginal stability and, employing the wall-crossing formula, extrapolate the spectrum to other regions of the moduli space. In the weak-coupling region, we recover the complete semiclassical spectrum derived in [3], thus confirming our starting assumption. The conclusions of a forthcoming analysis of the spectrum of the corresponding two-dimensional theory [11] are fully consistent with our results.

From our analysis, the following general picture emerges. For a given magnetic charge, there is a (“primary”) wall separating the strong-coupling region from the rest of the moduli space. Outside this wall, the spectrum expands and includes an infinite (“primary”) tower of dyons as well as quarks and W bosons. In addition, we show that if a particular condition on the complex masses is satisfied, there is one extra (“secondary”) tower of bound states consisting of a dyon and one or more quarks. Unlike the primary case, the wall-crossing formula shows that all the states in the extra tower cannot be created at a single wall. Rather, for every bound state in the tower, there is a corresponding (“secondary”) wall. We also show that each secondary wall separates the primary wall from the weak-coupling region, so that all walls must be traversed in passing between strong and weak coupling.

A particular configuration of \mathbb{Z}_n -symmetric masses, when all n masses form a regular polygon in the complex plane, can be analysed more explicitly: we find that there exists one secondary tower of bound states with one quark for odd n and no bound states for even n .

Let us introduce our conventions in the four-dimensional theory: $N_f = n + \tilde{n}$ is the total number of flavours, \vec{q}_e and \vec{q}_m are the vectors of electric and magnetic charges with n components (counted by I), \vec{S} is the vector of flavour charges with $n + \tilde{n}$ components (counted by i). The central charge is given by

$$Z_{(\vec{q}_e, \vec{q}_m, \vec{M})} = \vec{a} \vec{q}_e + \vec{a}_D \vec{q}_m + \vec{S} \vec{M} = \sum_{I=0}^{n-1} (a^I q_{eI} + a_{DI} q_{mI}^I) + \sum_{i=0}^{N_f-1} S_i M_i \quad (1)$$

where \vec{a} is the vacuum expectation value, \vec{a}_D is its magnetic dual, \vec{M} is the vector of flavour masses. We divide \vec{S} and \vec{M} into two pieces: \vec{s} and \vec{m} contain the first n components corres-

¹ By this, we mean that there is a single multiplet of $\mathcal{N} = (2, 2)$ SUSY in the two-dimensional theory associated with each massive multiplet of $\mathcal{N} = 2$ SUSY in the four-dimensional theory. With an appropriate identification of parameters, the masses of corresponding states agree exactly. For more details see [2, 3].

ponding to the massless quarks at the root of the Higgs branch, \vec{s} and \vec{m} contain the remaining \tilde{n} components; we distinguish the remaining \tilde{n} flavour components by putting a tilde above their masses, charges, and indices. The root of the Higgs branch is determined by setting $\vec{a} = -\vec{m}$; analogously, we define a magnetic dual mass $\vec{m}_D = -\vec{a}_D(\vec{a} = -\vec{m})$. Therefore, the central charge (1) reduces to

$$Z_{(\vec{\gamma}_e, \vec{\gamma}_m, \vec{s})} = \vec{m}\vec{\gamma}_e + \vec{m}_D\vec{\gamma}_m + \vec{s}\vec{m} = \sum_{I=0}^{n-1} (m^I \gamma_{eI} + m_{DI} \gamma_{mI}) + \sum_{\tilde{i}=0}^{\tilde{n}-1} s_{\tilde{i}} m_{\tilde{i}}, \quad (2)$$

$$\vec{\gamma}_e = -\vec{q}_e + \vec{s}, \quad \vec{\gamma}_m = -\vec{q}_m.$$

Now, for each BPS state, the complete set of (electric, magnetic, and flavour) charges is $\gamma = (\vec{\gamma}_e, \vec{\gamma}_m, \vec{s})$; if $\vec{s} = \vec{0}$, we will omit the third entry.

Our approach is mainly based on the Kontsevich–Soibelman wall-crossing formula [4]. For a given charge $\gamma = (\vec{\gamma}_e, \vec{\gamma}_m)$, we define the Kontsevich–Soibelman operator acting on the so-called Darboux coordinates \mathcal{X}_β (for any charge β) as

$$\mathcal{K}_\gamma : \mathcal{X}_\beta \rightarrow \mathcal{X}_\beta (1 - \sigma(\gamma) \mathcal{X}_\gamma)^{\langle \beta, \gamma \rangle} \quad (3)$$

where for any pair of charges, $\alpha = (\vec{\alpha}_e, \vec{\alpha}_m)$ and $\beta = (\vec{\beta}_e, \vec{\beta}_m)$, the symplectic product is defined as

$$\langle (\vec{\alpha}_e, \vec{\alpha}_m), (\vec{\beta}_e, \vec{\beta}_m) \rangle = -\vec{\alpha}_e \vec{\beta}_m + \vec{\alpha}_m \vec{\beta}_e, \quad (4)$$

and the quadratic refinement is defined as

$$\sigma(\gamma) = (-1)^{\vec{\gamma}_e \vec{\gamma}_m}. \quad (5)$$

Let $\Gamma(\vec{M})$ be the set of BPS states, depending on the set of masses \vec{M} . We associate the following operator to each point \vec{M} in the moduli space:

$$S = \prod_{\gamma \in \Gamma(\vec{M})} \mathcal{K}_\gamma^{\Omega(\gamma, \vec{M})} = \text{const} \quad (6)$$

where $\Omega(\gamma, \vec{M})$ is the degeneracy of the BPS state with charge γ ; all operators (i.e., their BPS rays) are ordered clockwise (equivalently, their central charges as complex vectors are ordered counterclockwise). The statement of the wall-crossing formula is that, although the spectrum and the ordering of operators change across the moduli space, the resulting product S is constant. In principle, knowing the set of charges on one side of the wall of marginal stability, we can compute them on the other side of the wall [5]. In practice, this proceeds via the use of known identities, such as (14) and (17) below, which apply for specific values of the symplectic product of the two states whose central charges become aligned at the wall.

2 The wall at strong coupling

Our starting point is the strong-coupling spectrum of the theory. As mentioned above, we start by assuming the spectrum implied by the 2d/4d conjecture of [2, 3]. For the moment, we consider only the BPS states corresponding to kinks interpolating between two neighbouring

vacua in the two-dimensional theory. Without loss of generality, we can set the magnetic charge to be equal to $(-1, 1, 0, 0, \dots)$. There are exactly n such states (plus charge conjugates). Electric charges are determined only up to a fixed shift [3]; for our purposes, it is convenient to choose this shift so that the charges of the states are

$$\begin{aligned}\pm\gamma_1 &= \pm((1, 0, 0, 0, \dots), (-1, 1, 0, 0, 0, \dots)), \\ \pm\gamma_2 &= \pm((0, 1, 0, 0, \dots), (-1, 1, 0, 0, 0, \dots)), \\ \pm\gamma_3 &= \pm((0, 0, 1, 0, 0, \dots), (-1, 1, 0, 0, 0, \dots)), \\ &\dots\end{aligned}\tag{7}$$

The moduli space includes Argyres–Douglas points [16] located on the boundary of the strong-coupling region. One of these states becomes massless at each of these points.

Consider the case of \mathbb{Z}_n -symmetric masses. We set

$$m_I = m_0 \exp \frac{2\pi i I}{n}, \tag{8}$$

where m_0 remains as a free parameter which interpolates between strong and weak coupling. Then, all central charges and walls of marginal stability can be determined as functions of m_0 , and the masses automatically obey

$$\sum_{I=0}^{n-1} m_I = 0. \tag{9}$$

A convenient expression for the magnetic dual masses in the \mathbb{Z}_n -symmetric case is [10]²

$$m_{DI} = e^{2\pi i I/n} \left(n \sqrt[n]{m_0^n + \Lambda^n} + \sum_{k=0}^{n-1} m_0 e^{2\pi i k/n} \log \frac{\sqrt[n]{m_0^n + \Lambda^n} - m_0 e^{2\pi i k/n}}{\Lambda} \right) \tag{10}$$

where the branch is fixed by requiring that for $x \in \mathbb{R}_+$, $\sqrt[n]{x} \in \mathbb{R}_+$ and $\log x \in \mathbb{R}$, as in [10]. Then, the Argyres–Douglas points, where the strong-coupling states (7) become massless, are located at [10]

$$m_0 = \Lambda \exp \frac{\pi i (2j+1)}{n}, \quad j \in \mathbb{Z} \tag{11}$$

(figure 1). As explained in [17], for \mathbb{Z}_n -symmetric masses, it is sufficient to consider m_0 belonging to the sector between two neighbouring Argyres–Douglas points, $\Lambda e^{\pi i/n}$ and $\Lambda e^{-\pi i/n}$, where γ_1 and γ_2 from (7) are massless.

Let us find out how the spectrum changes when \vec{M} crosses the primary wall of marginal stability, where the central charges of the first two dyons in (7), γ_1 and γ_2 , become aligned:

$$\frac{Z_{\gamma_1}}{Z_{\gamma_2}} \in \mathbb{R}_+ \tag{12}$$

(figure 2). Using the wall-crossing formula, we can compute the spectrum on the external side

² Although the analysis of [10] is for the corresponding two-dimensional theory, the resulting formula for the central charge is identical to that of the four-dimensional theory at the Higgs branch root [2].

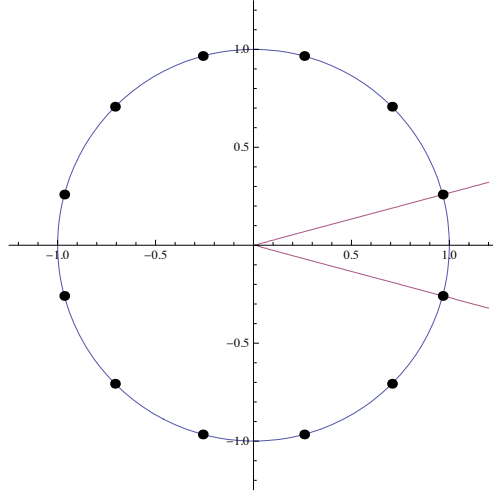


Figure 1: Argyres–Douglas points for \mathbb{Z}_{12} -symmetric masses in the m_0 plane.

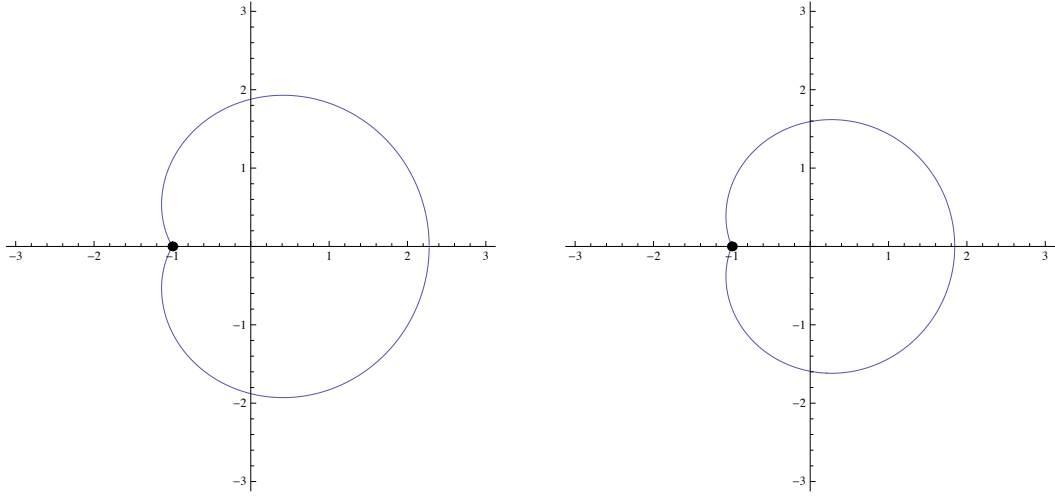


Figure 2: The primary walls of marginal stability for \mathbb{Z}_2 - and \mathbb{Z}_3 -symmetric masses [18, 10].

of the wall. The symplectic product is

$$\langle \gamma_1, \gamma_2 \rangle = -2, \quad (13)$$

hence, the relevant wall-crossing formula is essentially the same as the formula relating the strong- and weak-coupling spectra of the pure $SU(2)$ theory in four dimensions³ [4]:

$$\mathcal{K}_{-\gamma_2} \mathcal{K}_{\gamma_1} = \mathcal{K}_{\gamma_1} \mathcal{K}_{2\gamma_1 - \gamma_2} \mathcal{K}_{3\gamma_1 - 2\gamma_2} \mathcal{K}_{4\gamma_1 - 3\gamma_2} \dots \mathcal{K}_{\gamma_1 - \gamma_2}^{-2} \dots \mathcal{K}_{3\gamma_1 - 4\gamma_2} \mathcal{K}_{2\gamma_1 - 3\gamma_2} \mathcal{K}_{\gamma_1 - 2\gamma_2} \mathcal{K}_{-\gamma_2}. \quad (14)$$

In our notations, the part of the wall-crossing formula that changes across the wall takes the following form:

$$\begin{aligned} \mathcal{K}_{-(0,1),(-1,1)} \mathcal{K}_{((1,0),(-1,1))} &= \mathcal{K}_{((1,0),(-1,1))} \mathcal{K}_{((2,-1),(-1,1))} \mathcal{K}_{((3,-2),(-1,1))} \mathcal{K}_{((4,-3),(-1,1))} \\ &\dots \mathcal{K}_{((-1,1),(0,0))}^{-2} \dots \mathcal{K}_{-((-3,4),(-1,1))} \mathcal{K}_{-((-2,3),(-1,1))} \mathcal{K}_{-((-1,2),(-1,1))} \mathcal{K}_{-((0,1),(-1,1))} \end{aligned} \quad (15)$$

³ This reflects the fact that all the states involved have charges contained in an $SU(2)$ subgroup of the gauge group.

where we display only the first two components of electric and magnetic charges, the others being equal to zero. This relation shows that the spectrum outside the wall consists of a tower of dyons and a finite number of quarks and W bosons with charges

$$\begin{aligned} & \pm ((-\nu + 1, \nu, 0, 0, \dots), (-1, 1, 0, 0, \dots)), \\ & \pm ((-1, 1, 0, 0, \dots), (0, 0, 0, 0, \dots)). \end{aligned} \quad (16)$$

3 Bound states

In fact, the complete BPS spectrum is not limited to the primary tower of states found above: depending on the values of masses, there can also be secondary towers of bound states formed by a dyon and p quarks [3]. Creation (or, conversely, destruction) of these extra states is described by the pentagon formula:

$$\mathcal{K}_{\gamma_1} \mathcal{K}_{\gamma_2} = \mathcal{K}_{\gamma_2} \mathcal{K}_{\gamma_1 + \gamma_2} \mathcal{K}_{\gamma_1}, \quad \forall \langle \gamma_1, \gamma_2 \rangle = \pm 1 \quad (17)$$

where the new state $\gamma_1 + \gamma_2$ is created from γ_1 and γ_2 where one of the initial states is a quark, and the other one is either a dyon or a bound state consisting of a dyon and $p - 1$ quarks. This process occurs when

$$\frac{Z_{\gamma_1}}{Z_{\gamma_2}} \in \mathbb{R}_+. \quad (18)$$

We will find the resulting secondary walls and prove that they are located outside the primary wall and have to be crossed as the VEV moves from strong to weak coupling. The restriction on the wedge-product of the two interacting states in (17) allows us to determine which states can combine to form a bound state if the corresponding secondary wall exists.

Starting with the states constructed in the previous section, when $\tilde{n} = 0$, we can see that there can be two possible types of creation processes, both leading to the same set of new states:

$$\begin{aligned} 1 : & \quad ((-\nu + 1, \nu, 0, 0, 0, \dots), (-1, 1, 0, 0, 0, \dots)) + ((-1, 0, 1, 0, 0, \dots), (0, 0, 0, 0, 0, \dots)) \\ & \quad \leftrightarrow ((-\nu, \nu, 1, 0, 0, \dots), (-1, 1, 0, 0, 0, \dots)), \end{aligned} \quad (19)$$

$$\begin{aligned} 2 : & \quad ((-\nu, \nu + 1, 0, 0, 0, \dots), (-1, 1, 0, 0, 0, \dots)) + ((0, -1, 1, 0, 0, \dots), (0, 0, 0, 0, 0, \dots)) \\ & \quad \leftrightarrow ((-\nu, \nu, 1, 0, 0, \dots), (-1, 1, 0, 0, 0, \dots)). \end{aligned} \quad (20)$$

These are the bound states formed by a dyon and one quark. Explicitly, the walls of marginal stability (18) for these processes are determined by

$$1 : \quad \frac{-m_0 + m_2}{(-\nu + 1)m_0 + \nu m_1 - m_{D0} + m_{D1}} \in \mathbb{R}_+, \quad (21)$$

$$2 : \quad \frac{-m_1 + m_2}{-\nu m_0 + (\nu + 1)m_1 - m_{D0} + m_{D1}} \in \mathbb{R}_+. \quad (22)$$

For general N_f , there are additional bound states involving the remaining \tilde{n} flavours:

$$\begin{aligned} & ((-\nu + 1, \nu, 0, 0, 0, \dots), (-1, 1, 0, 0, 0, \dots)) + ((-1, 0, 0, 0, 0, \dots), (0, 0, 0, 0, 0, \dots), (1, 0, 0, \dots)) \\ & \quad \leftrightarrow ((-\nu, \nu, 0, 0, 0, \dots), (-1, 1, 0, 0, 0, \dots), (1, 0, 0, \dots)), \end{aligned} \quad (23)$$

$$\begin{aligned}
& ((-\nu, \nu + 1, 0, 0, 0, \dots), (-1, 1, 0, 0, 0, \dots)) + ((0, -1, 0, 0, 0, 0, \dots), (0, 0, 0, 0, 0, \dots), (1, 0, 0, \dots)) \\
& \leftrightarrow ((-\nu, \nu, 0, 0, 0, 0, \dots), (-1, 1, 0, 0, 0, \dots), (1, 0, 0, \dots)).
\end{aligned} \tag{24}$$

They are completely analogous to the ones above: the walls of marginal stability for these processes can be obtained by changing m_2 to \tilde{m}_0 in the previous formulae.

As has been discussed above, there can also be bound states formed by a dyon and p quarks:

$$\begin{aligned}
& ((-\nu + 1 + p, \nu, j_3, j_4, j_5, \dots), (-1, 1, 0, 0, 0, \dots), (\tilde{j}_1, \tilde{j}_2, \tilde{j}_3, \dots)), \\
& j_i(j_i - 1) = \tilde{j}_{\tilde{i}}(\tilde{j}_{\tilde{i}} - 1) = 0, \quad p + \sum_{i=2}^{n-1} j_i + \sum_{\tilde{i}=0}^{\tilde{n}-1} \tilde{j}_{\tilde{i}} = 0.
\end{aligned} \tag{25}$$

They exist if starting with the strong-coupling spectrum and moving into the weak-coupling region, $|p|$ different secondary walls of marginal stability (18) are crossed.

We need to find out if the processes (19) and (20), which we rewrite as

$$1 : \quad d_1 + q_1 \leftrightarrow b, \tag{26}$$

$$2 : \quad d_2 + q_2 \leftrightarrow b, \tag{27}$$

actually take place when the masses move from strong to weak coupling: to do this, we should check whether the secondary walls (18) are crossed, i.e., if the following conditions are satisfied somewhere outside the primary wall of marginal stability:

$$1 : \quad \arg Z_{d_1} = \arg Z_{q_1}, \tag{28}$$

$$2 : \quad \arg Z_{d_2} = \arg Z_{q_2}. \tag{29}$$

Note that Z_{q_j} ($j = 1$ or $j = 2$) is independent of the region in the moduli space, and $\arg Z_{d_j}$ changes continuously between the primary wall and the weak-coupling region, therefore, (28) (with $j = 1$) and (29) (with $j = 2$) are satisfied somewhere if in the complex plane, Z_{q_j} lies between the values of Z_{d_j} at the primary wall and in the weak-coupling limit.

To check if this is the case, it is convenient to start at the Argyres–Douglas point where either γ_1 or γ_2 in (7) becomes massless. Consider (28) first. For $\nu > 0$ and for $\nu \leq 0$, we start at the points s_1 and s_2 where γ_1 and γ_2 in (7) are massless, respectively. Near these points, the corresponding central charges of dyons can be approximated as

$$(Z_{d_1})_{s_1} \simeq \nu(-m_0 + m_1), \quad (Z_{d_1})_{s_2} \simeq (\nu - 1)(-m_0 + m_1). \tag{30}$$

Then, we continuously move the masses into the semiclassical region, where

$$(Z_{d_1})_w \simeq i(-m_0 + m_1). \tag{31}$$

From (30) and (31), we have

$$\lim_{g_{\text{eff}} \rightarrow 0} \arg \frac{(Z_{d_1})_w}{(Z_{d_1})_{s_l}} = (-1)^{l-1} \frac{\pi}{2}, \quad \left| \arg \frac{(Z_{d_1})_w}{(Z_{d_1})_{s_l}} \right| < \frac{\pi}{2}, \tag{32}$$

where the inequality is strict for any g_{eff} because the central charge receives corrections from its electric components at weak coupling. We can consider (29) analogously: (19) and (20) are related by changing $\nu \rightarrow -\nu$ and swapping $\gamma_{e0} \leftrightarrow \gamma_{e1}$. Comparing Z_{q_j} with (30) and (31) and using (32), we conclude that the walls (21, 22) exist when

$$\begin{aligned} 1: \quad \nu > 0: \quad \arg \frac{m_k - m_0}{m_1 - m_0} &\in \left(0, \frac{\pi}{2}\right), \\ \nu \leq 0: \quad \arg \frac{m_k - m_0}{m_1 - m_0} &\in \left(\frac{\pi}{2}, \pi\right), \end{aligned} \quad (33)$$

$$\begin{aligned} 2: \quad \nu \geq 0: \quad \arg \frac{m_k - m_1}{m_1 - m_0} &\in \left(0, \frac{\pi}{2}\right), \\ \nu < 0: \quad \arg \frac{m_k - m_1}{m_1 - m_0} &\in \left(\frac{\pi}{2}, \pi\right). \end{aligned} \quad (34)$$

Since the bound states with $\nu \neq 0$ do not exist at strong coupling, they appear at weak coupling if exactly one of the two walls (28, 29) is crossed. The states with $\nu = 0$, which exist at strong coupling, appear at weak coupling if either none or both walls (28, 29) are crossed. For all three cases, $\nu > 0$, $\nu < 0$, and $\nu = 0$, this means that the bound states in (19, 20) exist semiclassically when the following condition is satisfied:

$$0 < \text{Re} \frac{m_k - m_0}{m_1 - m_0} < 1 \quad (35)$$

(again, the inequality is strict because of (32)). This is precisely the semiclassical constraint derived in [3] from first principles. Analogously, if (35) holds for p different indices k and \tilde{k} , there are towers of bound states with p quarks (25) having $j_k = 1$ and $\tilde{j}_{\tilde{k}} = 1$ for these indices and $j_i = 0$ and $\tilde{j}_{\tilde{i}} = 0$ for all other i and \tilde{i} , in accordance with [3].

Applying this result to \mathbb{Z}_n -symmetric masses, it is easy to find the bound states in the weak-coupling limit. The constraint (35) reduces to

$$0 < \text{Re} \frac{e^{2\pi k i/n} - 1}{e^{2\pi i/n} - 1} = \text{Re} \frac{e^{2\pi(k-1/2)i/n} - e^{-\pi i/n}}{2i \sin(\pi/n)} < 1 \quad \Longleftrightarrow \quad -1 < \frac{\sin \frac{2\pi k - \pi}{n}}{\sin \frac{\pi}{n}} < 1. \quad (36)$$

Here, $(2\pi k - \pi)/n$ is a multiple of π/n , therefore, the inequality holds only for $(2\pi k - \pi)/n = \pi$, that is, for $k = (n+1)/2$. For \mathbb{Z}_{2l+1} -symmetric masses with $l \in \mathbb{N}$, this means that only the bound states formed by one quark with $\gamma_{e(l+1)} = 1$ are present (figure 3); for \mathbb{Z}_{2l} -symmetric masses, there are no bound states. We can go back to equations (33, 34) to find out which secondary walls of marginal stability exist in the case of \mathbb{Z}_{2l+1} -symmetric masses: (19) is realised for $\nu > 0$, (20) is realised for $\nu < 0$, and the corresponding walls are determined by (21) with $\nu > 0$ and (22) with $\nu < 0$ (plotted for \mathbb{Z}_3 in figure 4); in addition, all bound states with $\nu = 0$ not belonging to the tower of bound states decay between strong- and weak-coupling, as discussed above.

The coils corresponding to $\nu = \pm(p+1)$ and to $\nu = \pm p$ for $p \neq 0$ are consecutive sections of the same spiral. The clockwise and the counterclockwise spirals contain the coils with $\nu > 0$ and $\nu < 0$, respectively. This follows from the fact that the rotation by 2π in the moduli space changes the electric charges by 1.

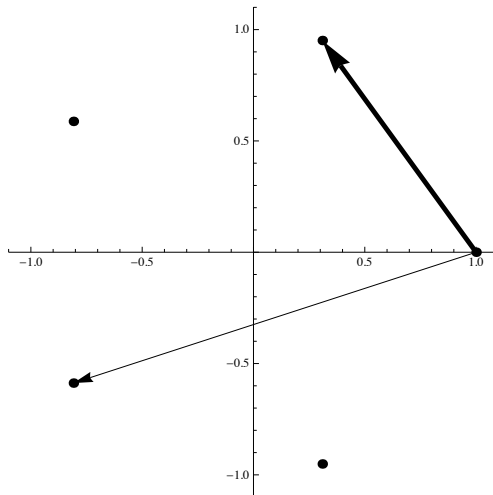


Figure 3: A dyon (thick vector connecting m_0 and m_1 , representing its central charge near the massless point at strong coupling for $\nu > 0$) and the quark that can form bound states with it (thin vector connecting m_0 and m_3 , equal to its central charge) in the case of \mathbb{Z}_5 -symmetric masses.

The wall-crossing formula discussed above places strong constraints on the presence of any extra BPS states and their possible decay processes. For example, in the four-dimensional model, additional states not belonging to the secondary tower with quantum numbers

$$((-\nu, \nu, 0, 0, \dots, 0, 1, 0, 0, \dots), (-1, 1, 0, 0, 0, \dots)), \quad (37)$$

analogous to the extra towers of state in the two-dimensional model discussed in [10], are not present. In particular, the wall-crossing formula certainly forbids the obvious simultaneous decay process for the tower of such states into the known quarks and dyons of the model:

$$((-\nu, \nu, 1, 0, 0, \dots), (-1, 1, 0, 0, 0, \dots)) \stackrel{?}{\leftrightarrow} ((0, 0, 1, 0, 0, \dots), (-1, 1, 0, 0, 0, \dots)) + \nu((-1, 1, 0, 0, 0, \dots), (0, 0, 0, 0, 0, \dots)). \quad (38)$$

To see this, suppose that the spectrum on right-hand side is correct, and we cross the wall in the other direction. The symplectic product of the two charges on the right-hand side is two, which means that the formula (14) should apply. One can then check, however, that this leads to a different decay process into states having magnetic charges greater than one (in one $SU(2)$ subgroup), which are certainly absent from the model.

In this paper, we do not consider the corresponding two-dimensional theory directly⁴. However, our spectrum in the 4d theory for \mathbb{Z}_n -symmetric masses coincides⁵ with the relevant 2d spectrum obtained in the forthcoming paper [11], and this agreement provides further support for the 2d/4d correspondence of [2, 3].

⁴ Note that a similar wall-crossing formula is believed to hold for 2d models of this type [19].

⁵ More precisely, the two spectra coincide up to minor discrepancies which originate in the precise assignment of charges to the particles in the strong coupling region.

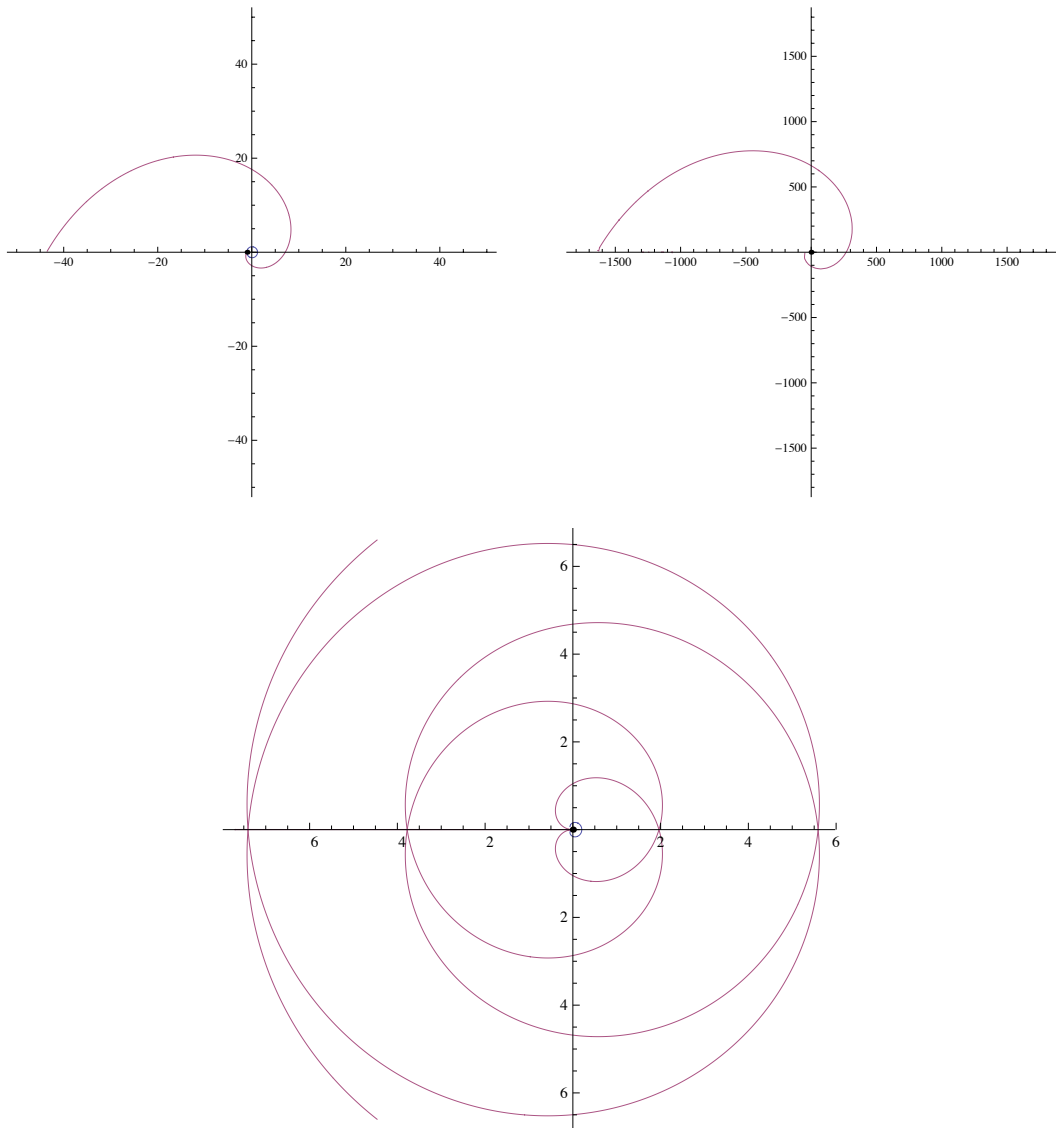


Figure 4: First row: secondary walls of marginal stability for \mathbb{Z}_3 -symmetric masses in the $m_0^3/|m_0|^2$ plane for $\nu = 1$ and $\nu = 2$ (the walls for $\nu = -1$ and $\nu = -2$ are their reflections across the real axis); second row: the two spirals and the primary wall with all radii scaled as $r \rightarrow \log r$. The plots correspond to the sector in figure 1.

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